Holographic Cosmological Constant and Dark Energy

Chao-Jun Feng

Institute of Theoretical Physics, Academia Sinica

P.O. Box 2735, Beijing 100080, China

and

Interdisciplinary Center of Theoretical Studies

USTC, Hefei, Anhui 230026, China

fengcj@itp.ac.cn

A general holographic relation between UV and IR cutoff of an effective field theory is proposed. Taking the IR cutoff relevant to the dark energy as the Hubble scale, we find that the cosmological constant is highly suppressed by a numerical factor and the fine tuning problem seems alleviative. We also use different IR cutoffs to study the case in which the universe is composed of matter and dark energy.

Sept. 2007

1. Introduction

Why cosmological constant observed today is so much smaller than the Planck scale? This is one of the most important problems in modern physics. In history, Einstein first introduced the cosmological constant in his famous field equation to achieve a static universe in 1917. After the discovery of the Hubble's law, the cosmological constant was no longer needed because the universe is expanding. Nowadays, the accelerating cosmic expansion first inferred from the observations of distant type Ia supernovae [1][2] has strongly confirmed by some other independent observations, such as the cosmic microwave background radiation (CMBR) [3] and Sloan Digital Sky Survey (SDSS) [4], and the cosmological constant returns back as a simplest candidate to explain the acceleration of the universe in 1990's.

In particle physics, the cosmological constant naturally arises as an energy density of the vacuum , which is evaluated by the sum of zero-point energies of quantum fields with mass m as follows

$$\rho_\Lambda = \frac{1}{2} \int_0^\Lambda \frac{4\pi k^2 dk}{(2\pi)^3} \sqrt{k^2 + m^2} \approx \frac{\Lambda^4}{16\pi^2}, \tag{1.1}$$

where $\Lambda \gg m$ is the UV cutoff. Usually the quantum field theory is considered to be valid just below the Planck scale: $M_p \sim 10^{18} \text{GeV}$, where we used deduced Planck mass $M_p^{-2} = 8\pi G$ for convenience. If we pick up $\Lambda = M_p$, we find that the energy density of the vacuum in this case is estimated as 10^{70}GeV^4, which is about 10^{117} orders of magnitude larger than the observation value 10^{-47}GeV^4. One may try to cancel it by introducing counter terms, however, this requires a fine tuning to adjust the energy density of the vacuum to the present energy density of the universe (for a classic review see [5], for a recent nice review see [6], and for a recent discussion see [7][8]). It seems that the number of the independent degrees of freedom of the quantum fields should not be very large [9][10].

Holographic principle [11] regards black holes as the maximally entropic objects of a given region and postulates that the maximum entropy inside this region behaves non-extensively, growing only as its surface area. Hence the number of independent degrees of freedom is bounded by the surface area in Planck units, so an effective field theory with UV cutoff Λ in a box with size L will make sense if it satisfies the Bekenstein entropy bound [12]

$$(L\Lambda)^3 \leq S_{BH} = \pi L^2 M_{pl}^2, \tag{1.2}$$

where $M_{pl}^{-2} \equiv G$ is the Planck mass and S_{BH} is the entropy of a black hole of radius L which acts as an IR cutoff. Cohen and collaborators [13] suggested that the total energy in a region of size L should not exceed the mass of a black hole of the same size

$$L^3 \Lambda^4 \leq L M_p^2 \,, \tag{1.3}$$

which can be simply rewritten as

$$(L\Lambda)^4 \leq L^2 M_p^2 \,. \tag{1.4}$$

This bound is much more stringent than the bound (1.2): when equation (1.4) is near saturation, the entropy of the quantum field is

$$S_{max} \approx S_{BH}^{3/4} \,. \tag{1.5}$$

Since we have limited knowledge about the holographic principle and we have not even know whether the holographic principle is right or not because we have only a few examples to realize it. The only successful example to my knowledge is the AdS/CFT correspondence. Mostly, one believes the holographic principle is right because it does not conflict with any observations so far. As a result we can not claim whether the bounds mentioned above as a consequence of the holographic principle is correct or not, and it may be too stringent or too loose due to some unknown reasons or some underlying theory. In section 2 we postulate a general bound which provides a mechanism to derive a very small vacuum energy from the principle of holography.

When the matter presents in the universe, the evolution of the dark energy (in this note we shall use terms the cosmological constant and the dark energy exchangeably) is sensitive to the chosen of the IR cutoff. When we take the event horizon as the IR cutoff, the result is very similar to the case studied in ref. [14] up to some corrections. As long as the vacuum dominates the energy density in the later time, it should be small as we discussed in section 2.

This paper is organized as follows. In section 2 we postulate a general relation between the UV an IR cutoff and the smallness of the cosmological constant shall be explained. In the next section we use three different IR cutoffs to study the evolution and the equation of state of the dark energy. In the final section we will give some discussions.

2. General bound and the cosmological constant

In this note we postulate a general relation of UV and IR cutoff as follows

$$(L\Lambda)^n \leq L^2 M_p^2 \,, \tag{2.1}$$

where n is a dimensionless parameter that comes from some underlying theory. When $n = 3, 4$ equation (2.1) is reduced to (1.2) and (1.4) respectively, and we shall see that the final consistent n is slightly deviation from 4 but without any fine tuning. Of course we can not say anything about this unknown theory yet, since we do not even know whether there really exists such a theory or not, but we shall see that if the relation (2.1) is correct it will provide a mechanism to derive a small cosmological constant without any fine tuning. There are some works trying to solve the cosmological constant problem from the holographic principle, for instance see [15][16][14][17][18][19], but they do not consider the general case of the relation (2.1) between UV and IR cutoff.

The largest Λ allowed here is the one saturating the inequality (2.1):

$$\Lambda = L^{\frac{2}{n}-1} M_p^{\frac{2}{n}} \,. \tag{2.2}$$

Then the energy density of the vacuum $\rho_\Lambda \sim \Lambda^4$ is

$$\rho_\Lambda = 3c^2 L^{\frac{8}{n}-4} M_p^{\frac{8}{n}} \,, \tag{2.3}$$

where a numerical constant $3c^2$ is introduced in the above equation for convenience. From (1.1) one can see the value of c^2 is naturally neither very large nor very small. Of course there is also a constant in the equation (2.2), but naturally such a constant could not be very large or very small and it will not affect the final conclusion since one can absorb this constant into c^2.

The dynamics of the universe is described by the Einstein field equations. The observations indicate that the universe is homogeneity and isotropy on large scales and the generic metric respecting these symmetries is the Friedmann- Robertson-Walker (FRW) metric given by

$$ds^2 = -dt^2 + a^2(t) \left[\frac{dr^2}{1 - Kr^2} + r^2(d\theta^2 + sin^2\theta d\phi^2) \right] \,, \tag{2.4}$$

where $a(t)$ is scale factor with cosmic time t. The coordinates r, θ and ϕ are known as comoving coordinates. A freely moving particle comes to rest in these coordinates. The

constant K in the metric (2.4) describes the geometry of the spatial section of space time, and $K = +1, 0, -1$ corresponds to closed, flat and open universe respectively.

Consider an ideal perfect fluid with energy density ρ and pressure p as the source of the energy momentum tensor and solve the Einstein equation with the metric (2.4), we find the famous Friedmann equations

$$H^2 = \frac{\rho}{3M_p^2} - \frac{K}{a^2} \tag{2.5}$$

and

$$\dot{H} = -\frac{\rho + p}{2M_p^2} + \frac{K}{a^2} \tag{2.6}$$

where $H \equiv \dot{a}/a$ is the Hubble constant. Since observations indicate the universe is flat, i.e. the critical energy density is almost equal to 1, we will only consider the flat case $K = 0$ in the following. In fact (2.6) can be derived with the help of the continuity equations respecting the conservation of the energy momentum as the consequence of the Bianchi identities:

$$\dot{\rho} + 3H(\rho + p) = 0. \tag{2.7}$$

To indicate the application of the relation (2.1), we consider such a situation in which we assumed that the IR cutoff is the Hubble scale H^{-1} and the vacuum dominates the universe:

$$3M_p^2 H^2 = 3c^2 H^{4-\frac{8}{n}} M_p^{\frac{8}{n}}, \tag{2.8}$$

which can be easily solved

$$\left(\frac{H}{M_p}\right)^{2-\frac{8}{n}} = \frac{1}{c^2}. \tag{2.9}$$

For a given n ($n \neq 4$) H is a constant, so the universe is de Sitter space. While $n = 4$, then $c^2 = 1$ and H could be any value since $n = 4$ is an unstable point, thus we can get very small value of cosmological constant when n is slightly deviation from 4, and there is not any fine tuning problem in such a difference, we shall see this in the following. We would like to emphasize that n and c here are determined by some underlying reasons as a normal number, by normal number we mean the number is not very large like 10^{10} or small like 10^{-10}, so there is no need for us to adjust them to produced a small vacuum energy. If our postulation is right, it should make it. Usually c^2 is not equal to 1 , so the energy density of the vacuum is
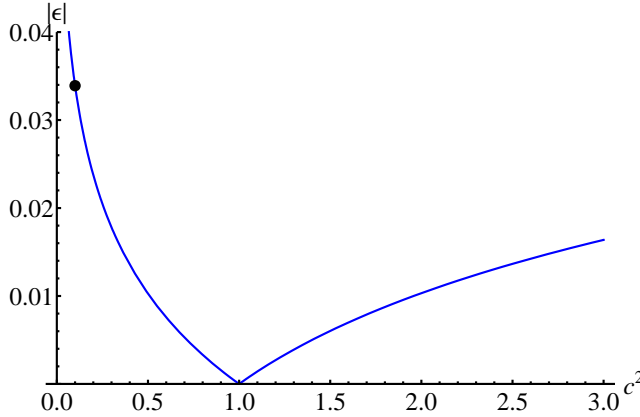
$$\rho_\Lambda = 3c^2 H^{4-\frac{8}{n}} M_p^{\frac{8}{n}} = 3(c^2)^{\frac{n}{4-n}} M_p^4 \tag{2.10}$$

4

where we have used (2.9). If $c^2 < 1$ and $n < 4$, the energy density can be highly suppressed and much smaller than the Planck scale energy density M_p^4 in the limit of $n \to 4$. If $c^2 > 1$ and $n > 4$, one can also get a very small energy density. In the other case the value of energy density is ruled out by the observations because it is too large. To illustrate that n has the slight but not fine-tuning difference from 4, we give a concrete example in the following.

Assuming $c^2 = 0.1$, and if

$$n = n_{0.1} = \frac{4}{1 - \frac{\ln c^2}{117 \ln 10}} \approx 3.966 \,, \tag{2.11}$$

the energy density is roughly $10^{-117} M_p^4$. Let $n = 4 + \epsilon$, one can see $\epsilon = -0.03$ in this example. In fact, such a difference $|\epsilon|$ would not be a extreme small number, namely, a number like 10^{-10}, as long as c^2 is not very closed to 1, so there is no fine tuning problem here, one can see this property in Figure. 1.

Figure 1. *The difference $|\epsilon| \equiv |n - 4|$ vs. c^2. The value of c^2 and ϵ on the curve will deduce a small energy density, namely, $10^{-117} M_p^4$. The black point on the curve is the point $c^2 = 0.1$, $|\epsilon| \approx 0.03$.*

At first glance it seems that there is a fine tuning problem here: for a given c^2, one should adjust the value of n to be very closed to $n_{0.1}$ for example. But it is not the case, because we do not need to adjust c^2 and n in fact, the resulting energy density is a consequence or a prediction of the theory rather than a input. The figure above only indicates that the underlying theory should contain a constant number n whose value is near 4 without any fine tuning.

Let's have a look at equation (2.1)

$$(L\Lambda)^{4+\epsilon} \leq L^2 M_p^2 \,. \tag{2.12}$$

It seems that we are living in the fractal dimension spacetime rather than 4, if one regards the power of $L\Lambda$ as the dimension of the world we living due to the simple fact that the max entropy in 4 dimension spacetime is roughly $L^3\Lambda^3$ and in 3 dimension spacetime is $L^2\Lambda^2$. It is amazing if this explanation is correct because it is so counterintuitive. But we are forced to reconsider the stability of the orbits of planets like the earth because there are no stable solutions to keep the earth rounding the sun, in other words the Inverse Square Law is not hold in high dimensions. Things is also bad in low dimensional worlds because no one can live in 2 dimension space. Problems may be disappeared here because of the fractal dimension. Another possibility is that in the following.

Take the correction term to the R. H. S. of (2.12) as

$$L^3\Lambda^4 \leq (L\Lambda)^{-\epsilon} L M_p^2. \tag{2.13}$$

The factor $(L\Lambda)^{-\epsilon}$ may come from the some unknown theory. Since this factor has something to do with the cosmological constant which can be considered as a consequence of quantum gravity [20], one could guess this factor may come from the correction of quantum gravity theory. Maybe these two possibilities are the same thing, but there is no evidence here now.

3. Dark energy with the presence of matter

With matter present, the Friedmann equation reads

$$3M_p^2 H^2 = \rho_m + \rho_\Lambda \,, \tag{3.1}$$

where ρ_m, the energy density of matter, satisfies the continuity equation:

$$\ln' \rho_m + 3 = 0 \,. \tag{3.2}$$

where prime denotes the derivative with respect to $\ln a$. The dimensionless energy density of matter is defined as $\Omega_m = \rho_m/(3M_p^2 H^2)$, so it satisfies the following equation

$$\ln' \Omega_m = -3 - 2\ln' H \,, \tag{3.3}$$

where we have used the continuity equation (3.2).

Define a dimensionless energy density

$$\Omega_\Lambda \equiv \frac{\rho_\Lambda}{3M_p^2 H^2} = c^2 (LM_p)^{\frac{8}{n}-4} H^{-2} M_p^2. \tag{3.4}$$

Then the Friedmann equation (3.1) is

$$\Omega_m + \Omega_\Lambda = 1. \tag{3.5}$$

Derivative the logarithm of (3.4) with respect to $\ln a$ as follows

$$\ln' \Omega_\Lambda = -2 \left[(2 - \frac{4}{n}) \ln' L + \ln' H \right] \approx -2 \left[\ln'(LH) + \frac{\epsilon}{4} \ln' L \right] \tag{3.6}$$

From (3.3), (3.5) and (3.6) we can derive a equation of H and L

$$\ln' H + (1 + \frac{\epsilon}{4}) \Omega_\Lambda \ln' L + \frac{3}{2}(1 - \Omega_\Lambda) = 0 \tag{3.7}$$

and once we know the another relation of L and H, we can solve this equation.

Measuring w as in $\rho_\Lambda \sim a^{-3(1+w)}$, we have the index w given by

$$w = -1 - \frac{1}{3} \left(\frac{d \ln \rho_\Lambda}{d \ln a} + \frac{1}{2} \frac{d^2 \ln \rho_\Lambda}{d(\ln a)^2} \ln a \right), \tag{3.8}$$

up to the second order and the derivatives are taken at the present time $a_0 = 1$. From (2.3) we find

$$\frac{d \ln \rho_\Lambda}{d \ln a} = \frac{\dot\rho_\Lambda}{H\rho_\Lambda} = \left(\frac{8}{n} - 4 \right) \frac{\dot L}{LH} \approx -2(1 + \frac{\epsilon}{4}) \ln' L \tag{3.9}$$

thus we get the ration of pressure to energy density w as

$$w \approx -1 + \frac{2}{3}(1 + \frac{\epsilon}{4}) \ln' L. \tag{3.10}$$

For the acceleration of the universe $w < -1/3$, the R. H. S. of (3.10) should satisfy

$$(1 + \frac{\epsilon}{4}) \ln' L < 1. \tag{3.11}$$

and for a increasing of Ω_Λ the R.H.S. of (3.6) should be positive

$$\ln'(LH) + \frac{\epsilon}{4} \ln' L < 0. \tag{3.12}$$

From (3.11) one can see that, when the IR cutoff L is smaller than 1, the universe is acceleration definitely and this vacuum energy will eventually dominate the universe. This happens when we regards the event horizon as a natural cutoff as the IR cutoff, and Miao's work in [14] has already indicated such a character.

In the following we will study the property of the vacuum energy with three different IR cutoffs : Hubble scale H^{-1}, particle horizon R_p and event horizon R_h, since these IR cutoffs naturally arise when one studies the universe. The definition of R_p and R_h is given by

$$R_p(t) = a(t) \int_0^t \frac{dt'}{a(t')} \tag{3.13}$$

$$R_h(t) = a(t) \int_t^\infty \frac{dt'}{a(t')} \tag{3.14}$$

3.1. Case1: $L = H^{-1}$

In this case, the vacuum energy behaves almost like the matter, which means it's equation of state is very similar to that of the matter up to some corrections, and we find this correction will lead to an evolution of w, but it will take a long time for w to be -1. In other words, the vacuum energy will not accelerate the universe until it almost completely dominates the universe. Since the calculation discussed above is straightforward, we simply give the final result as follows.

From (3.6) and (3.7) the equation of the energy density is found to be

$$\frac{\Omega'_\Lambda}{\Omega_\Lambda} = -\frac{3\epsilon}{4} \left[\frac{1 - \Omega_\Lambda}{1 - (1 + \frac{\epsilon}{4})\Omega_\Lambda} \right], \tag{3.15}$$

where $L = H^{-1}$ was used. When $\epsilon < 0$ the dimensionless energy density of the vacuum is increasing with time. This equation can be solved easily as

$$\ln \Omega_\Lambda + \frac{\epsilon}{4} \ln(1 - \Omega_\Lambda) = -\frac{3\epsilon}{4} \ln a + x_0 \tag{3.16}$$

If we set $a_0 = 1$ at the present time, x_0 is equal to the L.H.S. of (3.16) with Ω_Λ replaced by Ω_Λ^0 , namely, $x_0 = \ln \Omega_\Lambda^0 + \frac{\epsilon}{4} \ln(1 - \Omega_\Lambda^0)$. As time draws by, Ω_Λ increases to 1, the second term on the L.H.S. of (3.16) is the important term, we find, for large a

$$\Omega_\Lambda = 1 - e^{4x_0/\epsilon} a^{-3}. \tag{3.17}$$

Since the universe is dominated by the dark energy for large a, we have

$$\rho_\Lambda \sim \rho_c = \frac{\rho_m}{1 - \Omega_\Lambda} = \frac{\rho_m^0 a^{-3}}{1 - \Omega_\Lambda} \tag{3.18}$$

Thus, using (3.17) in the above relation

$$\rho_\Lambda = e^{-4x_0/\epsilon} \rho_m^0 \tag{3.19}$$

which is too large compared with the observation value of Ω_Λ^0, if we require a acceleration universe, namely $\epsilon < 0$.

For small a, matter dominates, the important term on the L.H.S. of (3.16) is the first term, we find

$$\Omega_\Lambda = a^{-3\epsilon/4} e^{x_0} \tag{3.20}$$

thus

$$\rho_\Lambda = \Omega_\Lambda \rho_c = \Omega_\Lambda \rho_m = e^{x_0} \rho_m^0 a^{-3(1+\epsilon/4)}, \tag{3.21}$$

here ϵ is much smaller than 1, so the evolution of the vacuum energy is roughly a^{-3} the same as the matter when a is small. In other words w is almost zero when matter presents.

Up to the second order, the equation of sate is described by

$$w = \frac{\epsilon}{4 - (4 + \epsilon)\Omega_\Lambda^0} + \frac{3\epsilon^2(1 + \epsilon/4)\Omega_\Lambda^0(1 - \Omega_\Lambda^0)}{32(1 - (1 + \epsilon/4)\Omega_\Lambda^0)^3} z \tag{3.22}$$

where we used $\ln a = -\ln(1 + z) \sim -z$.

Specifying to the case $\epsilon = -0.03$ and plugging the optional value $\Omega_\Lambda^0 = 0.73$ into (3.22),

$$w = -0.027 + 7.9 \times 10^{-4} z. \tag{3.23}$$

It seems that the Hubble scale is not a suitable IR cutoff.

3.2. Case2: $L = R_p$

If we take the particle horizon as the IR cutoff, the situation is not much changed from the Hubble scale case. Since

$$\ln' L = \ln' R_p = 1 + \frac{1}{R_p H} \tag{3.24}$$

9

the equation of state from (3.10) is

$$w = -\frac{1}{3} + \frac{\epsilon}{6} + \frac{2(1 + \frac{\epsilon}{4})}{3R_p H} \tag{3.25}$$

which is larger than $(-1/3 + \epsilon/6)$. It seems that the universe is hardly to accelerate in this case.

From (3.6) and (3.7) we find

$$\ln' \Omega_\Lambda = -2(1 - \Omega_\Lambda) \left[(1 + \frac{\epsilon}{4})(1 + \frac{1}{R_p H}) - \frac{3}{2} \right], \tag{3.26}$$

so the vacuum energy will dominates at later time if

$$R_p H > 2 \left(1 + \frac{3\epsilon}{4} \right) \tag{3.27}$$

It seems that the particle horizon is not a suitable IR cutoff either.

3.3. Case3: $L = R_h$

In this case the situation is changed, namely we can get an accelerating universe as follows and firstly one can simply see that

$$\ln' L = \ln' R_h = 1 - \frac{1}{R_h H}, \tag{3.28}$$

where the minus sign in (3.28) is the main difference from (3.24) , so the equation of state (3.25) is changed to be

$$w = -\frac{1}{3} + \frac{\epsilon}{6} - \frac{2(1 + \frac{\epsilon}{4})}{3R_h H} \tag{3.29}$$

which is smaller than $(-1/3 + \epsilon/6)$. It seems that the universe is able to accelerate in this case. Here the term $\epsilon/6$ will slightly change the value of w, and this is a correction to that in [14].

From (3.6) and (3.7) we find

$$\ln' \Omega_\Lambda = -2(1 - \Omega_\Lambda) \left[(1 + \frac{\epsilon}{4})(1 - \frac{1}{R_h H}) - \frac{3}{2} \right], \tag{3.30}$$

so the vacuum energy will dominates if

$$\frac{1}{R_p H} > -\frac{1}{2} \left(1 - \frac{3\epsilon}{4} \right) \tag{3.31}$$

10

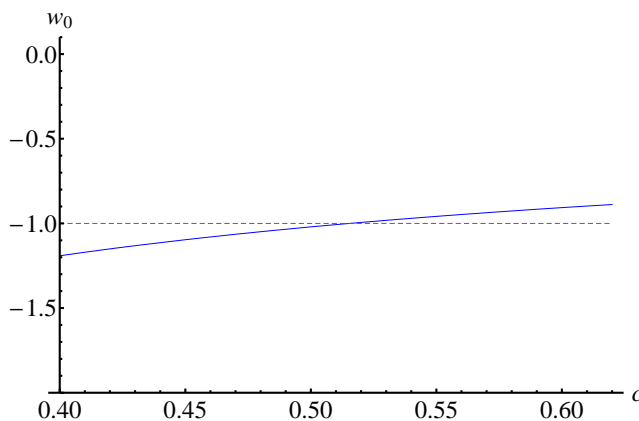which is always hold. Use the definition of Ω_Λ in (3.4), we find (3.30) becomes

$$\ln{}' \Omega_\Lambda = -2(1 - \Omega_\Lambda) \left[(1 + \frac{\epsilon}{4})(1 - \frac{\sqrt{\Omega_\Lambda}}{c}(R_h M_p)^{\epsilon/4}) - \frac{3}{2} \right], \tag{3.32}$$

which can not be solved analytically. The approximate solution when ϵ is small will reduce to the result in [14]. The corresponding equation of state from (3.29) will be

$$w = -\frac{1}{3} + \frac{\epsilon}{6} - \frac{2(1 + \frac{\epsilon}{4})}{3c} \sqrt{\Omega_\Lambda} \, (R_h M_p)^{\epsilon/4} . \tag{3.33}$$

Where we have used (3.4). If $\epsilon = 0$, the above equation is the same as that in [14], so terms containing ϵ are corrections with slight effects.

When the vacuum completely dominates the universe at last, the universe is a de Sitter space and the event horizon is roughly the inverse of the Hubble constant at that time t_0, namely $R_h \sim H_0^{-1}$. If we take today's value of the Hubble constant $H_0^4 \sim 10^{-117} M_p^4$, the factor $(R_h M_p)^{\epsilon/4}$ in (3.33) is roughly ~ 0.60 where we have used $\epsilon \sim -0.03$. Taking the present value of $\Omega_\Lambda^0 = 0.73$ and $c = 0.5$ then $w_0 \approx -1.02$. Figure. 2 shows the relation between w_0 and c, but notice that here c as a constant is from some unknown theory rather than adjusted.

Figure 2. *The dash line denotes the line $w_0 = -1$. The solid line indicates w_0 vs. c when $\epsilon = -0.03$ and the factor $(R_h M_p)^{\epsilon/4} \sim 0.6$.*

4. Discussions

The application of holographic principle discussed in the present paper alleviates the cosmological constant problem. When the vacuum dominates the universe, the energy

density could be very small due to the number of n and c^2 come from some underlying reasons. This provide a mechanism to explain why the cosmological constant is so small. In other words, one can get a very small energy density consistent with the observation value by this mechanism. We give an example and argue that there is no fine tuning problem in this mechanism. It should be emphasized that n and c^2 are given numbers rather than adjusted.

When matter presents as a component in the universe, this vacuum energy play a role as the dark energy. The evolution of the dark matter here is sensitive to the IR cutoff. We have used three different cutoff and find the result is consistent with [14], namely the event horizon is a suitable IR cutoff for the energy density to accelerate the universe.

At first glance it seems that if $\ln' L$ is non-positive the universe is definitely accelerating from (3.11). However, if $\ln' L < 0$, it means there is a shrinking IR cutoff, so the cutoff will be smaller and smaller as time draws by and we can see less and less stars and galaxies. This is absurd. If $\ln' L = 0$, it means there is a universal IR cutoff of the universe, by universal we mean the cutoff is independent of cosmic time, then the energy density of the dark matter is a constant, namely $w = -1$, but there is no evidence that we have such a universal IR cutoff. In a word, $\ln' L$ should be positive.

Acknowledgements

References

[1] A. G. Riess *et al.* [Supernova Search Team Collaboration], Astron. J. **116**, 1009 (1998) [arXiv:astro-ph/9805201].

[2] S. Perlmutter *et al.* [Supernova Cosmology Project Collaboration], Astrophys. J. **517**, 565 (1999) [arXiv:astro-ph/9812133].

[3] D. N. Spergel *et al.* [WMAP Collaboration], Astrophys. J. Suppl. **170**, 377 (2007) [arXiv:astro-ph/0603449].

[4] For a recent date release see, J. K. Adelman-McCarthy *et al.* [SDSS Collaboration], arXiv:0707.3413 [astro-ph].

[5] S. Weinberg, Rev. Mod. Phys. **61**, 1 (1989).

[6] S. M. Carroll, Living Rev. Rel. **4**, 1 (2001) [arXiv:astro-ph/0004075].

[7] J. Polchinski, arXiv:hep-th/0603249.

[8] R. Bousso, arXiv:0708.4231 [hep-th].

[9] G. 't Hooft, arXiv:gr-qc/9310026.

[10] L. Susskind, J. Math. Phys. **36**, 6377 (1995) [arXiv:hep-th/9409089].

[11] For a review of holographic principle see, R. Bousso, Rev. Mod. Phys. **74**, 825 (2002) [arXiv:hep-th/0203101].

[12] J. D. Bekenstein, Phys. Rev. D **7**, 2333 (1973); J. D. Bekenstein, Phys. Rev. D **23**, 287 (1981).

[13] A. G. Cohen, D. B. Kaplan and A. E. Nelson, Phys. Rev. Lett. **82**, 4971 (1999) [arXiv:hep-th/9803132].

[14] M. Li, Phys. Lett. B **603**, 1 (2004) [arXiv:hep-th/0403127]. Q. G. Huang and M. Li, JCAP **0503**, 001 (2005) [arXiv:hep-th/0410095]. Q. G. Huang and M. Li, JCAP **0408**, 013 (2004) [arXiv:astro-ph/0404229]. B. Chen, M. Li and Y. Wang, Nucl. Phys. B **774**, 256 (2007) [arXiv:astro-ph/0611623]. J. f. Zhang, X. Zhang and H. y. Liu, arXiv:0708.3121 [hep-th].

[15] P. Horava and D. Minic, Phys. Rev. Lett. **85**, 1610 (2000) [arXiv:hep-th/0001145].

[16] Y. S. Myung, Phys. Lett. B **578**, 7 (2004) [arXiv:hep-th/0306180].

[17] R. Horvat, Phys. Rev. D **70**, 087301 (2004) [arXiv:astro-ph/0404204].

[18] M. M. Sheikh-Jabbari, Phys. Lett. B **642**, 119 (2006) [arXiv:hep-th/0605110].

[19] D. Podolsky and K. Enqvist, arXiv:0704.0144 [hep-th].

[20] E. Witten, arXiv:hep-ph/0002297.